

Gamow-Teller transitions and proton-neutron pair correlation in $N = Z = \text{odd } p\text{-shell nuclei}$

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Abstract

We have studied the Gamow-Teller (GT) transitions from $N = Z + 2$ neighbors to $N = Z = \text{odd}$ nuclei in p -shell region by using isospin-projected and $\beta\gamma$ -constraint antisymmetrized molecular dynamics combined with generator coordinate method. The calculated GT transition strengths from 0^+1 states to 1^+0 states such as ${}^6\text{He}(0_1^+1) \rightarrow {}^6\text{Li}(1_1^+0)$, ${}^{10}\text{Be}(0_1^+1) \rightarrow {}^{10}\text{B}(1_1^+0)$, and ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_2^+0)$ exhaust more than 50% of the sum rule. These $N = Z + 2$ initial states and $N = Z = \text{odd}$ final states are found to dominantly have $S = 0, T = 1$ nn pairs and $S = 1, T = 0$ pn pairs, respectively. Based on two-nucleon (NN) pair picture, we can understand the concentration of the GT strengths as the spin-isospin-flip transition $nn(S = 0, T = 1) \rightarrow pn(S = 1, T = 0)$ in LS -coupling scheme. The GT transition can be a good probe to identify the spin-isospin partner states with nn pairs and pn pairs of $N = Z + 2$ and $N = Z = \text{odd}$ nuclei, respectively.

I. INTRODUCTION

Proton and neutron (pn) correlation is one of the key phenomena to understand properties of $N = Z = \text{odd}$ nuclei (see Ref. [1] and references therein). In particular, a deuteron-like $T = 0$ pn pair plays an important role in low-lying states of light $N = Z = \text{odd}$ nuclei. Recently, a three-body model calculation of two nucleons with a doubly magic core nucleus has been performed to study low-lying states of $N = Z = \text{odd}$ nuclei, and the result indicates that deuteron-like $S = 1, T = 0$ and di-neutron-type $S = 0, T = 1$ pairs (LS -coupling pn pairs) are predominantly formed at the surface of double magic cores such as ^{16}O [2]. Moreover, in our previous work, we have studied pn correlation in ^{10}B and found the low-lying $T = 0$ and $T = 1$ states dominantly have the $S = 1, T = 0$ and $S = 0, T = 1$ pairs around a 2α core, respectively [3]. It indicates that LS -coupling scheme is better than jj -coupling scheme to understand pn pairs in light $N = Z = \text{odd}$ nuclei even though the LS -coupling pn pairs may change to jj -coupling pn pairs in heavy-mass systems because of the spin-orbit mean potential.

Gamow-Teller transition is one of the useful observables to verify the LS -coupling pn pairs because it is sensitive to the spin-isospin configuration. For light $N = Z = \text{odd}$ nuclei, the GT operator flips nucleon spins and isospins of a pair and changes the di-neutron-type nn pair to the deuteron-like pn pair. This type of GT transitions correspond to the collectivity of the proton-neutron pair and dominate the GT sum rule if the core parts are spin-isospin saturated systems and give no contribution to the GT transition. These modes are different from so-called Gamow-Teller giant resonances which are contributed by collectivity of the excess neutrons. Recently, the super-allowed GT transitions in the low-energy region have been observed and discussed in relation to the pn correlations [4]. The LS -coupling pn pairs may play an important role to the low-energy super-allowed GT transitions. However, there are a few theoretical works to systematically investigate LS -coupling pn pairs in light $N = Z = \text{odd}$ nuclei, though proton-neutron pairing correlations in medium and heavy mass $N = Z = \text{odd}$ nuclei have been discussed in jj -coupling scheme with mean field approaches [5–9].

The authors have a constructed new framework, isospin-projected $\beta\gamma$ -constraint anti-symmetrized molecular dynamics ($T\beta\gamma$ -AMD), which is useful in description of a pn pair in deformed or clustered systems [3]. In this paper, we investigate the GT transitions and

pn pairs in p -shell nuclei applying $T\beta\gamma$ -AMD to ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$. We discuss strong GT transitions in terms of NN pair in LS -coupling scheme and propose an interpretation of the initial and final states as spin-isospin partners. A particular attention is paid on the role of non-zero intrinsic spin ($S = 1$) of the $T = 0$ pn pair and its coupling with the orbital angular momentum of pn center of mass motion and that of core rotation.

The paper is organized as follows. We briefly explain our framework in Sect. II. We show the results of nuclear properties of energy spectra, $B(M1)$, $B(E2)$ and $B(\text{GT})$ in Sect. III. We discuss the strong GT transitions in terms of NN pair in LS -coupling scheme in Sect. IV by analyzing the obtained wavefunctions. A summary and an outlook are given in Sect. V.

II. METHOD

A. $T\beta\gamma$ -AMD

For $N = Z = \text{odd}$ nuclei, we apply $T\beta\gamma$ -AMD [3] in order to deal with the pn pair formation as well as nuclear deformation and clustering. For $N = Z + 2$ nuclei, we use $\beta\gamma$ -constraint AMD [10], which has been used for structure studies of light neutron-rich nuclei as well as $Z = N = \text{even}$ nuclei. We here briefly explain the formulation of $T\beta\gamma$ -AMD. Details of two methods, $T\beta\gamma$ -AMD and $\beta\gamma$ -AMD are described in Refs. [3, 10].

In the original framework of AMD, a basis wavefunction is written by a Slater determinant of Gaussian wave packets,

$$|\Phi(\beta, \gamma)\rangle = \mathcal{A}[|\phi_1\rangle |\phi_2\rangle \cdots |\phi_A\rangle], \quad (1)$$

$$|\phi_i\rangle = \left(\frac{2\nu}{\pi}\right)^{\frac{3}{4}} \exp\left[-\nu\left(i - \frac{Z_i}{\sqrt{\nu}}\right)^2\right] |\xi_i\rangle |\tau_i\rangle. \quad (2)$$

In the present work, we use $\nu = 0.235$ for p -shell nuclei as used in Refs. [10–15]. In $T\beta\gamma$ -AMD, we perform parity and isospin (${}^\pi T$) projections before variation as

$$|\Phi^{\pi T}(\beta, \gamma)\rangle = \hat{P}^\pi \hat{P}^T |\Phi(\beta, \gamma)\rangle, \quad (3)$$

where \hat{P}^π and \hat{P}^T are parity projection operator and isospin projection operator, respectively. For the ${}^\pi T$ -projected AMD wavefunction, we perform energy variation under the constraint on quadrupole deformation parameters $\beta\gamma$ and obtain the optimum solution for each set of β and γ values. In order to obtain wavefunctions for the n th $J^\pi T$ state (denoted by $J_n^\pi T$),

we superpose the angular momentum eigenstates projected from the obtained wavefunctions $|\Phi^{\pi T}(\beta_i, \gamma_i)\rangle$,

$$|J_n^{\pi T}; M\rangle = \sum_{iK} c_n^{iK} \hat{P}_{MK}^J |\Phi^{\pi T}(\beta_i, \gamma_i)\rangle, \quad (4)$$

where \hat{P}_{MK}^J is the angular momentum projection operator. Here, the parameters, β and γ , are treated as generator coordinates in the generator coordinate method (GCM), and the K -mixing is taken into account. We call this method $T\beta\gamma$ -AMD+GCM.

B. Effective interactions

We use the Hamiltonian

$$H = K - K_{\text{cm}} + V_c + V_{LS} + V_{\text{Coulomb}}, \quad (5)$$

where K is the kinetic energy, K_{cm} is the kinetic energy of the center of mass, and V_c , V_{LS} , and V_{Coulomb} are the central, spin-orbit, and Coulomb forces, respectively. For the central and spin-orbit forces, we use the effective nucleon-nucleon (NN) forces same as those used for ^{10}B in the previous work [3]. Namely, we use the Volkov No. 2 force of the central force with the Majorana exchange parameter $m = 0.6$ and the G3RS force of the spin-orbit force with the strength parameters $u_1 = -u_2 = 1300$ MeV.

For the Bartlett and Heisenberg parameters, b and h , of the Volkov No. 2 force, we use $b = h = 0.125$ for ^6Li , which reproduce the S -wave NN scattering lengths in the $T = 0$ and $T = 1$ channels. For ^{10}B and ^{14}N , we adopt a parameterization $b = h = 0.06$ phenomenologically modified so as to describe energy difference between the lowest $T = 0$ and $T = 1$ states in each nucleus. The parameters b and h control the ratio (f) of the central force in the $T = 0$ channel to that in the $T = 1$ channel. The present choices, $b = h = 0.125$ and $b = h = 0.06$, give the ratios $f = 1.67$ and 1.27 , respectively. The decrease of f is consistent with the naive expectation that the $T = 0$ interaction is somewhat suppressed by a nuclear medium effect. We should comment that, even though relative position between $T = 1$ and $T = 0$ spectra is sensitive to b and h , we obtain almost the same energy spectra in each isospin channel and also qualitatively similar results for structure properties of ^{10}B and ^{14}N in the cases of $b = h = 0.125$ and $b = h = 0.06$.

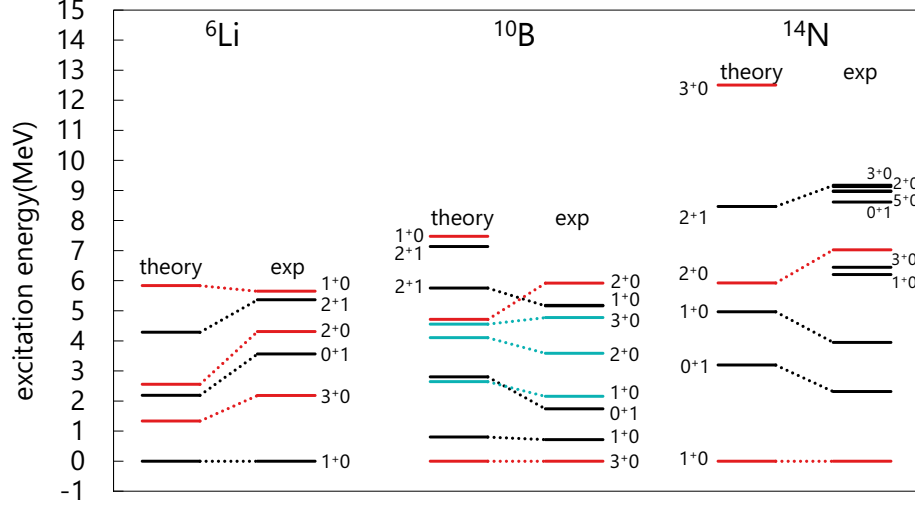


FIG. 1. Spectra of ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$ calculated by $T\beta\gamma\text{-AMD+GCM}$ and those of the experimental data [16–18].

III. RESULTS

Calculated energy spectra of ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$ obtained by $T\beta\gamma\text{-AMD+GCM}$ are shown in Fig. 1 compared with experimental spectra. The present calculation reasonably reproduces the low-energy spectra of these nuclei.

The calculated binding energies, magnetic dipole moments (μ), electric quadrupole moments (Q), and $E2$ and $M1$ transition strengths of ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$ are listed in Table I together with experimental data. The present calculation quantitatively or qualitatively reproduces the experimental data of these properties.

μ moments and $B(M1)$ as well as the GT transition strengths are observables that sensitively reflect spin configurations. The calculated μ moments of the ground states, ${}^6\text{Li}(1_1^+0)$, ${}^{10}\text{B}(3_1^+0)$, and ${}^{14}\text{N}(1_1^+0)$, and that of ${}^{10}\text{B}(1_1^+0)$ agree well to the experimental data. For the $M1$ transition in ${}^6\text{Li}$, the remarkably large $B(M1; 0_1^+1 \rightarrow 1_1^+0)$ is well reproduced by the calculation. For ${}^{10}\text{B}$, the $M1$ transitions between $T = 1$ and $T = 0$ states are qualitatively described although quantitative reproduction is not satisfactory in the present calculation. For ${}^{14}\text{N}$, the present calculation describes the general trend of the relatively strong $M1$ transitions for $1_2^+0 \rightarrow 0_1^+1$ and $2_1^+1 \rightarrow 2_1^+0$ compared with those for other transitions.

For ${}^{10}\text{B}$, the calculated Q moment and $B(E2; 3_2^+0 \rightarrow 1_1^+0)$ are large consistently with the experimental data because of the prolate deformation. The prolate deformation of ${}^{10}\text{B}$ is

caused by formation of a 2α core as shown later.

TABLE I: Binding energies, μ and Q moments, and $M1$ and $E2$ transition strengths of ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$. The calculated values obtained by $T\beta\gamma$ -AMD+GCM are shown. Experimental data are taken from [16–18].

Observable	$T\beta\gamma$ -AMD+GCM	Exp
${}^6\text{Li}$		
$ E(1_1^+0) $ (MeV)	29.55	31.99
$\mu(1_1^+0)$ (μ_N)	0.87	0.82
$Q(1_1^+0)$ ($e\text{ fm}^2$)	0.09	-0.08
$B(E2; 3_1^+0 \rightarrow 1_1^+0)$	3.79	10.69
$B(M1; 0_1^+1 \rightarrow 1_1^+0)$	13.73	15.43
$B(E2; 2_1^+0 \rightarrow 1_1^+0)$	5.15	4.40
$B(M1; 2_1^+1 \rightarrow 1_1^+0)$	0.01	0.15
${}^{10}\text{B}$		
$ E(3_1^+0) $ (MeV)	60.35	64.75
$\mu(3_1^+0)$ (μ_N)	1.83	1.80
$\mu(1_1^+0)$ (μ_N)	0.84	0.63
$Q(3_1^+0)$ ($e\text{ fm}^2$)	8.45	8.47
$B(E2; 1_1^+0 \rightarrow 3_1^+0)$	4.03	4.15
$B(M1; 0_1^+1 \rightarrow 1_1^+0)$	14.98	7.52
$B(M1; 1_2^+0 \rightarrow 0_1^+1)$	0.05	0.19
$B(E2; 1_2^+0 \rightarrow 1_1^+0)$	9.23	15.61
$B(E2; 1_2^+0 \rightarrow 3_1^+0)$	2.02	1.70
$B(E2; 2_1^+0 \rightarrow 3_1^+0)$	0.34	1.15
$B(E2; 3_2^+0 \rightarrow 1_1^+0)$	10.56	19.71
$B(M1; 2_1^+1 \rightarrow 2_1^+0)$	1.84	2.52
$B(M1; 2_1^+1 \rightarrow 1_2^+0)$	2.60	3.06
$B(M1; 2_1^+1 \rightarrow 1_1^+0)$	0.31	0.32

Observable	$T\beta\gamma$ -AMD+GCM	Exp
^{14}N		
$ E(1_1^+0) $ (MeV)	108.60	104.66
$\mu(1_1^+0)$ (μ_N)	0.34	0.40
$Q(1_1^+0)$ ($e\text{ fm}^2$)	0.53	1.93
$B(M1; 0_1^+1 \rightarrow 1_1^+0)$	0.76	0.05
$B(M1; 1_2^+0 \rightarrow 0_1^+1)$	3.72	1.79
$B(E2; 1_2^+0 \rightarrow 1_1^+0)$	3.25	4.41
$B(E2; 2_1^+0 \rightarrow 1_1^+0)$	2.95	3.61
$B(M1; 2_1^+1 \rightarrow 2_1^+0)$	4.65	1.74
$B(M1; 2_1^+1 \rightarrow 1_1^+0)$	0.00	0.59

In order to calculate GT transition strengths, we apply $\beta\gamma$ -AMD and obtain wavefunctions for the ground and excited states of ^6He , ^{10}Be , and ^{14}C , which are isobaric analogue states of $T = 1$ states of ^6Li , ^{10}B , and ^{14}N . Table II shows the calculated $B(\text{GT})$:

$$B(\text{GT}) = \frac{1}{2J_i + 1} \left| \left\langle J_f \left\| \sum_i \sigma^i \tau^i \right\| J_i \right\rangle \right|^2. \quad (6)$$

In the present paper, we define $B(\text{GT})$ by matrix elements of the spin and isospin operators without the factor $(g_A/g_V)^2$. For all low-lying states of ^6Li , ^{10}B , and ^{14}N , we find $T = 0$ states that have strong GT transitions with large percentages of the sum rule $\sum B(\text{GT}) = 3(N - Z) = 6$. These final states in $Z = N = \text{odd}$ nuclei can be regarded as “spin-isospin partners” of the corresponding $T = 1$ initial states because they are approximately spin-isospin-flipped states having spatial configurations similar to the initial states. The concept of the spin-isospin partners is an extension of isobaric analogue state (IAS) to the GT transition. The assignments of the spin-isospin partners in the following discussions are based on the calculated GT transition strengths and also spin and orbital configurations in LS -coupling scheme of NN pairs.

For ^6Li , the GT transition from the ground state $^6\text{He}(0_1^+1)$ to 1_1^+0 exhausts a large fraction of the sum rule consistently with the experimental data, whereas that to 1_2^+0 is weak. This fact indicates that the ground states of ^6Li and ^6He are almost ideal spin-isospin partners. For the GT transitions from the excited state, $^6\text{He}(2_1^+1)$, we obtain strong transitions to

1_2^+0 , 2_1^+0 , and 3_1^+0 . The summation of $B(\text{GT})$ values for these three states is about 50% of the sum rule, and therefore, these states are regarded as the set of spin-isospin partners with $J^\pi = \{1^+, 2^+, 3^+\}$ of the ${}^6\text{He}(2_1^+1)$.

Also for ${}^{10}\text{B}$, the GT transition strength from the ground state ${}^{10}\text{Be}(0_1^+1)$ is concentrated to 1_1^+0 . For the GT transitions from the excited states ${}^{10}\text{Be}(2_1^+1)$ and ${}^{10}\text{Be}(2_2^+1)$, significant strengths are obtained for transitions to 1_2^+0 , 1_3^+0 , 2_1^+0 , 2_2^+0 , 3_1^+0 , and 3_2^+0 , which can be assigned to two sets of spin-isospin partners with $J^\pi = \{1^+, 2^+, 3^+\}$ of ${}^{10}\text{Be}(2_1^+1)$ and ${}^{10}\text{Be}(2_2^+1)$. In particular, the transitions from ${}^{10}\text{Be}(2_1^+1)$ are significantly strong to 1_2^+0 and 3_2^+0 indicating that these states are regarded as spin-isospin partner states. In the transitions from ${}^{10}\text{Be}(2_2^+1)$, the strengths to 1_3^+0 and 3_1^+0 are significantly large, and hence, these states can be assigned to the spin-isospin partners of ${}^{10}\text{Be}(2_2^+1)$. For assignment of $J^\pi = 2^+$ states, the strengths $2_1^+1 \rightarrow 2_1^+0$, $2_1^+1 \rightarrow 2_2^+0$, $2_2^+1 \rightarrow 2_1^+0$, and $2_2^+1 \rightarrow 2_2^+0$ are comparable. It indicates that two $J^\pi = 2^+$ states corresponding to the partner states of ${}^{10}\text{Be}(2_1^+1)$ and ${}^{10}\text{Be}(2_2^+1)$ are strongly mixed with each other.

For ${}^{14}\text{N}$, the strongest GT transition from the ground state ${}^{14}\text{C}(0_1^+1)$ is obtained for 1_2^+0 consistently with the experimental data. It indicates that not 1_1^+0 but 1_2^+0 of ${}^{14}\text{N}$ is the spin-isospin partner in the $A = 14$ systems. Compared with the dominant transition $0_1^+1 \rightarrow 1_2^+0$, the calculated transition $0_1^+1 \rightarrow 1_1^+0$ is relatively minor but it does not reproduce the anomalously small value of the experimental datum. For the transitions from the first excited state 2_1^+1 , strengths to 1_1^+0 , 2_1^+0 , and 3_1^+0 are significant and the sum of them exhausts more than 80% of the sum rule; thus, these states are regarded as the spin-isospin partners.

TABLE II: Calculated values of $B(\text{GT})$ defined in Eq. (6) are shown. Experimental data are taken from [18–22]. The value in the parenthesis is the experimental datum for the mirror transition $^{10}\text{C}(0_1^+1) \rightarrow ^{10}\text{B}(1_1^+0)$.

Initial→Final	$T\beta\gamma\text{-AMD+GCM}$	Exp
$^6\text{He} \rightarrow ^6\text{Li}$		
$0_1^+1 \rightarrow 1_1^+0$	5.31	3.02
$0_1^+1 \rightarrow 1_2^+0$	0.00	–
$2_1^+1 \rightarrow 1_1^+0$	0.01	–
$2_1^+1 \rightarrow 3_1^+0$	0.97	–
$2_1^+1 \rightarrow 2_1^+0$	1.00	–
$2_1^+1 \rightarrow 1_2^+0$	1.10	–
$^{10}\text{Be} \rightarrow ^{10}\text{B}$		
$0_1^+1 \rightarrow 1_1^+0$	4.95	(2.20)
$0_1^+1 \rightarrow 1_2^+0$	0.15	–
$0_1^+1 \rightarrow 1_3^+0$	0.00	–
$2_1^+1 \rightarrow 3_1^+0$	0.63	0.07
$2_1^+1 \rightarrow 1_1^+0$	0.06	–
$2_1^+1 \rightarrow 1_2^+0$	0.81	–
$2_1^+1 \rightarrow 2_1^+0$	0.77	–
$2_1^+1 \rightarrow 3_2^+0$	1.71	–
$2_1^+1 \rightarrow 1_3^+0$	0.26	–
$2_1^+1 \rightarrow 2_2^+0$	0.86	–
$2_2^+1 \rightarrow 3_1^+0$	1.54	0.85
$2_2^+1 \rightarrow 1_1^+0$	0.01	–
$2_2^+1 \rightarrow 1_2^+0$	0.23	–
$2_2^+1 \rightarrow 2_1^+0$	0.71	–
$2_2^+1 \rightarrow 3_2^+0$	0.26	–
$2_2^+1 \rightarrow 1_3^+0$	0.82	–
$2_2^+1 \rightarrow 2_2^+0$	0.79	–

Initial→Final	$T\beta\gamma$ -AMD+GCM	Exp
$^{14}\text{C} \rightarrow ^{14}\text{N}$		
$0_1^+1 \rightarrow 1_1^+0$	0.30	3.64×10^{-6}
$0_1^+1 \rightarrow 1_2^+0$	4.32	1.70
$2_1^+1 \rightarrow 1_1^+0$	1.13	0.17
$2_1^+1 \rightarrow 2_1^+0$	1.77	–
$2_1^+1 \rightarrow 3_1^+0$	2.35	–

IV. DISCUSSION

In the previous section, we have discussed the assignments of spin-isospin partners focusing on the strong GT transitions in $A = 6$, $A = 10$, and $A = 14$ nuclei. In this section, we discuss detailed features of NN pairs in the spin-isospin-partner states.

A. Intrinsic structure and spatial distribution of a proton-neutron pair

In the obtained wavefunctions for the $A = 6$, $A = 10$, and $A = 14$ systems, NN pairs are found to be formed around α , 2α , and ^{12}C cores, respectively. In order to see the spatial distribution of the $S = 1, T = 0$ and $S = 0, T = 1$ NN pairs in the spin-isospin partners, we calculate two-particle density $\rho_{ST}(r)$ at the identical point r in the intrinsic states defined as

$$\rho_{ST}(r) = \frac{\langle \Phi^T(\beta, \gamma) | \hat{\rho}_{ST}(r) | \Phi^T(\beta, \gamma) \rangle}{\langle \Phi^T(\beta, \gamma) | \Phi^T(\beta, \gamma) \rangle}, \quad (7)$$

$$\hat{\rho}_{ST}(r) \equiv \sum_{ij} \hat{P}_{ij}^S \hat{P}_{ij}^T \delta(r - \hat{r}_i) \delta(r - \hat{r}_j), \quad (8)$$

where \hat{P}_{ij}^S and \hat{P}_{ij}^T are the spin and isospin projection operators for two particles. We define the two-nucleon-pair density $\rho_{NN}(r) \equiv \rho_{10}(r) - \rho_{01}(r)$ to cancel NN pair contributions from α clusters which contain the same numbers of $S = 1, T = 0$ NN pairs as those of $S = 0, T = 1$ NN pairs. With this definition, positive (negative) regions of $\rho_{NN}(r)$ indicate $S = 1, T = 0$ ($S = 0, T = 1$) pn -pair distributions in $T = 0$ ($T = 1$) states. In Fig. 2, we show the two-nucleon-pair density $\rho_{NN}(r)$ together with the one-body density distribution

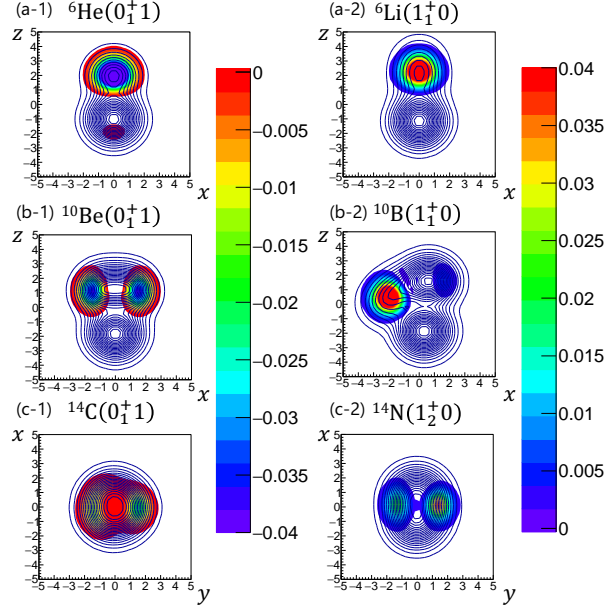


FIG. 2. The colored contours show two-nucleon-pair density $\rho_{NN}(r)$ of (a-1) ${}^6\text{He}(0_1^+1)$, (a-2) ${}^6\text{Li}(1_1^+0)$, (b-1) ${}^{10}\text{Be}(0_1^+1)$, (b-2) ${}^{10}\text{B}(1_1^+0)$, (c-1) ${}^{14}\text{C}(0_1^+1)$, and (c-2) ${}^{14}\text{N}(1_2^+0)$. The blue contours show the one-body density distribution $\rho(r)$.

in the single Slater-determinant state which has the largest overlap in the $\beta\gamma$ plane with the wavefunction for each of the ground 0_1^+1 states of ${}^6\text{He}$, ${}^{10}\text{Be}$, and ${}^{14}\text{C}$ and their spin-isospin partner 1^+0 states of the $N = Z = \text{odd}$ nuclei.

In the ground state of ${}^6\text{Li}$, an α particle and a $T = 0$ pn pair are formed as seen in Fig. 2(a-2). The pn pair spatially develops away from the α core and it shows deuteron-like nature. Also in ${}^6\text{He}$, the two-neutron pair appears around an α core (Fig. 2(a-1)). The spatial distribution of the nn pair density in ${}^6\text{He}$ is quite similar to that of the pn pair density in ${}^6\text{Li}$ indicating that these states are good spin-isospin partner states, in which the two-nucleon spin S and isospin T flip from $nn(T = 1, S = 0)$ in ${}^6\text{He}$ to $pn(T = 0, S = 1)$ in ${}^6\text{Li}$.

In ${}^{10}\text{B}(1_1^+0)$, two α clusters and a $T = 0$ pn pair are formed (see Fig. 2(b-2)). The $T = 0$ pn pair develops away from the 2α core similarly to ${}^6\text{Li}$, whereas the nn pair in ${}^{10}\text{Be}(0_1^+1)$ is not so developed spatially but is distributed at the nuclear surface showing a feature of two p -orbit neutrons (Fig. 2(b-1)). Although the single Slater-determinant state with the largest overlap for ${}^{10}\text{Be}(0_1^+1)$ shows less development of the two-nucleon pair than that for ${}^{10}\text{B}(1_1^+0)$, however, in the $\beta\gamma$ -AMD+GCM result, the spatially developed nn pair

components are largely mixed because the two-neutron pair can move away from the 2α core along a plateau toward a finite γ region in the $J^\pi = 0^+$ energy surface of ^{10}Be . As a result, the nn pair distribution in $^{10}\text{Be}(0_1^+1)$ has large overlap with the pn pair distribution in ^{10}B , and therefore, these states have the strong GT transition and are regarded as the partner states.

In $^{14}\text{N}(1_2^+0)$ (see Fig. 2(c-2)), a $T = 0$ pn pair is distributed at the surface of the oblatelly deformed ^{12}C core. In $^{14}\text{C}(0_1^+1)$, the nn pair density around the ^{12}C core shows distribution similar to the pn pair in $^{14}\text{N}(1_2^+0)$. The NN pairs in $^{14}\text{N}(1_2^+0)$ and $^{14}\text{C}(0_1^+1)$ show no spatial development and dominantly consist of p -orbit nucleons. If we consider a ^{16}O core, these states can be understood as two-hole pairs in the p -shell of the ^{16}O core.

Let us discuss spatial development of the NN pairs with A increasing in the $A = 6$, $A = 10$, and $A = 14$ systems. In the 0^+1 ground states of $N = Z + 2$ nuclei, the nn pair is mostly developed spatially in the $A = 6$ nucleus and comes down to the p -shell configurations in $A = 10$ and $A = 14$ nuclei with increase of the mass number. In the partner 1^+ states of the $N = Z = \text{odd}$ nuclei, the spatially developed $T = 0$ pn pair is prominent in the $A = 6$ nucleus and it more or less weakens but still remains even in the $A = 10$ nucleus, and finally comes down to the p -shell configuration in the $A = 14$ nucleus. This result reflects the feature that the $T = 0$ pn pairs in $N = Z = \text{odd}$ nuclei are robuster than nn pairs in $N = Z + 2$ nuclei. Indeed, the $T = 0$ pn pairs are described well by LS -coupling scheme, whereas nn pairs are somewhat broken from LS -coupling scheme and contain mixing of jj -coupling components, in particular, in the $A = 10$ and $A = 14$ nuclei as shown later in analysis of spin configurations.

B. pn pairs in LS -coupling scheme and spin-isospin partners

To quantitatively discuss the spin and orbital configurations, we show the expectation values of the squared intrinsic spin and orbital angular momentum ($\langle S^2 \rangle$ and $\langle L^2 \rangle$) in Table III. Note that $\langle S^2 \rangle$ approximately indicates the expectation value of the squared intrinsic spin of a NN pair around a core because core contribution is minor in the present case: the obtained states of the $A = 6$ and $A = 10$ nuclei are understood by two particles around $S = 0$ cores such as α and 2α and those of the $A = 14$ nuclei are approximately interpreted as two-hole states of ^{16}O . $T = 1$ states of the $N = Z = \text{odd}$ nuclei have almost same

expectation values as those of the $N = Z + 2$ nuclei because they are isobaric analogue states.

In the obtained states of the $A = 6$, $A = 10$, and $A = 14$ nuclei, the spin expectation values of $T = 0$ ($T = 1$) states are close to the value $\langle S^2 \rangle = 2$ ($\langle S^2 \rangle = 0$) for $S = 1$ ($S = 0$) component. It implies that LS -coupling NN pairs are formed as leading components in particular in light nuclei. As the mass number increases, the $T = 1, S = 0$ NN pairs in LS -coupling scheme are somewhat broken into jj -coupling pairs because of the spin-orbit mean potential. We can see this systematics especially in $\langle S^2 \rangle$ for the 0_1^+1 states. ${}^6\text{He}(0_1^+1)$ has almost pure $S = 0$ component with only 6% mixing of $S = 1$ component estimated from $\langle S^2 \rangle = 0.12$. However, ${}^{14}\text{C}(0_1^+1)$ has a broken $S = 0$ two-hole pair with significant $S = 1$ component up to 27%. In contrast to the $T = 1$ states, the LS -coupling pn pairs in the $T = 0$ states is not broken; the $S = 0$ mixing is found to be less than 6% for all the $T = 0$ states. This result implies that $T = 0, S = 1$ pn pairs are robuster than $T = 1, S = 0$ NN pairs. Even though NN pairs are not necessarily ideal LS -coupling pairs, they have LS -coupling features as major components and can be qualitatively understood by LS -coupling scheme.

For the orbital angular momentum, values of $\langle L^2 \rangle \approx 0$ and $\langle L^2 \rangle \approx 6$ indicate dominant $L = 0$ and $L = 2$ components, respectively. In $A = 6$ nuclei, the total orbital angular momentum L is contributed only by the orbital angular momentum L_{NN} of the NN pair because the α core is spherical. Therefore, the ground states ${}^6\text{He}(0_1^+1)$ and ${}^6\text{Li}(1_1^+0)$ are understood well by $S = 0, T = 1$ and $S = 1, T = 0$ NN pairs moving around the α in $L_{NN} = 0$ wave, whereas the excited states ${}^6\text{He}(2_1^+1)$ and ${}^6\text{Li}(1_2^+0, 2_1^+0, 3_1^+0)$ contain $S = 0, T = 1$ and $S = 1, T = 0$ NN pairs in $L_{NN} = 2$ wave. In the $A = 10$ nuclei, not only L_{NN} but also collective rotation of the 2α core with the orbital angular momentum L_{core} contributes to L . For ${}^{10}\text{Be}(0_1^+1)$ and ${}^{10}\text{B}(1_1^+0)$, $\langle L^2 \rangle \approx 0$ indicates that these states can be approximately described by the $S = 0, T = 1$ and $S = 1, T = 0$ NN pairs with $L_{NN} = L_{\text{core}} = 0$. The orbital angular momentum $L \approx 2$ of ${}^{10}\text{Be}(2_1^+1)$ mainly comes from the core rotation $L_{\text{core}} = 2$, whereas that of ${}^{10}\text{Be}(2_2^+1)$ is contributed mainly by $L_{NN} = 2$ from the NN pair rotation because the former and the latter states are a member of the $K = 0$ ground band and that of the $K = 2$ side band, respectively. It means that ${}^{10}\text{Be}(0_1^+1)$ and ${}^{10}\text{Be}(2_2^+1)$ are described by $S = 0$ nn pairs in $L_{NN} = 0$ and $L_{NN} = 2$ waves, respectively, and ${}^{10}\text{Be}(2_1^+1)$ is understood by a $S = 0$ nn pair with the rotating 2α core ($L_{\text{core}} = 2$).

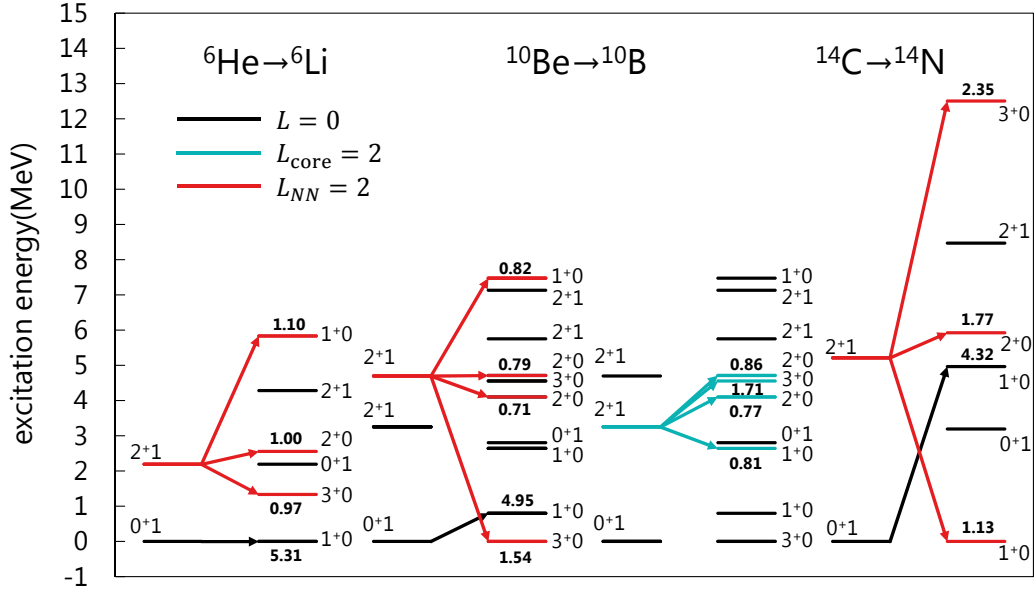


FIG. 3. GT transitions ${}^6\text{He} \rightarrow {}^6\text{Li}$, ${}^{10}\text{Be} \rightarrow {}^{10}\text{B}$, and ${}^{14}\text{C} \rightarrow {}^{14}\text{N}$ calculated by $T\beta\gamma\text{-AMD+GCM}$.

The corresponding spin-isospin partners in ${}^{10}\text{B}$ should have $S = 1, T = 0$ pn pairs with consistent spatial configurations. For the $A = 14$ systems, the dominant components of ${}^{14}\text{C}(0_1^+1)$ and ${}^{14}\text{N}(1_2^+1)$ have two holes in ${}^{16}\text{O}$ coupled to be $S = 0, T = 1$ and $S = 1, T = 0$ pairs in $L_{NN} = 0$ wave, whereas those of ${}^{14}\text{C}(2_1^+1)$ and ${}^{14}\text{N}(1_1^+0, 2_1^+0, 3_1^+0)$ are understood by $S = 0, T = 1$ and $S = 1, T = 0$ two-hole pairs in $L_{NN} = 2$ wave.

Based on LS -coupling scheme of NN pairs, we can easily understand spin-isospin partners and their strong GT transitions. The GT operator changes intrinsic spin configuration with $\Delta S = 1$ from $T = 1$ states to $T = 0$ states but it does not affect orbital configurations. In case of $L_{\text{core}} = 0$, nn pairs in $[L_{NN} = 0, S_{NN} = 0]_{J=0}$ initial states change directly into $T = 0$ pn pairs in $[L_{NN} = 0, S_{NN} = 1]_{J=1}$ states with strengths of the sum rule value: $\sum_n B(\text{GT}; 0^+1 \rightarrow 1_n^+0) = 6$ provided that core nuclei are spin-isospin saturated states and do not contribute to the GT transitions. Similarly, we can easily understand the spin-isospin partners of $[L_{NN} = 2, S_{NN} = 0]_{J=2}$ initial states and $[L_{NN} = 2, S_{NN} = 1]_{J=1,2,3}$ final states. Although J in the final states is not unique because of angular momentum coupling of $S_{NN} = 1$ with nonzero L_{NN} , we can again obtain the sum rule: $\sum_{J=1,2,3} \sum_n B(\text{GT}; 2^+1 \rightarrow J_n^+0) = 6$. It should be pointed out that, since $S = 1$ pn pairs in $L_{NN} = 2$ wave feel spin-orbit mean potentials from core nuclei, energy spectra of the final $J = 1, 2, 3$ states show spin-orbit splitting which plays an essential role to lower the $T = 0$ states into the ground

states in ^{10}B and ^{14}N . For deformed nuclei, we can also consider spin-isospin partners of $[L_{\text{core}} = 2, L_{NN} = 0, S_{NN} = 0]_{J=0}$ initial states and $[L_{\text{core}} = 2, L_{NN} = 0, S_{NN} = 1]_{J=1,2,3}$ final states.

If the NN pairs are broken into the jj -coupling pairs, the concentration of GT transition strengths does not occur because initial states change into various (jj') configurations. In other words, the concentration of GT transition strengths to specific final states is a good measure for realization of LS -coupling NN pairs. Based on the LS -coupling picture of NN pairs, we assigned the spin-isospin partners for $T = 1$ states and $T = 0$ states with strong GT transition strengths which are qualitatively characterized by $\Delta T = 0$, $\Delta S = 1$, $\Delta L = 0$ transitions.

Fig. 3 shows the calculated energy spectra and the GT transitions for the spin-isospin partners in $A = 6$, $A = 10$, and $A = 14$. The GT transition from $^6\text{He}(0_1^+1)$ to $^6\text{Li}(1_1^+0)$ is enhanced because these states have the same $L_{NN} = 0$ nature. For the excited states, the GT transitions from $^6\text{He}(2_1^+1)$ to $^6\text{Li}(1_2^+0, 2_1^+0, 3_1^+0)$ are strong because of the transition from the $T = 1, S = 0$ pair to the $T = 0, S = 1$ pair in the dominant $L_{NN} = 2$ component. The sum of the GT strengths from $^6\text{He}(2_1^+1)$ exhausts a large fraction of the sum rule value indicating that the nature of spin-isospin partners still remains also in the excited states. In the energy spectra of $^6\text{Li}(1_2^+0, 2_1^+0, 3_1^+0)$, the ordering of 3_1^+0 , 2_1^+0 , and 1_2^+0 is easily understood by the spin-orbit splitting for the $S = 1$ pn pairs in $L_{NN} = 2$.

We can also understand the GT transition from $^{10}\text{Be}(0_1^+1)$ to $^{10}\text{B}(1_1^+0)$ in the picture of LS -coupling NN pairs as GT transition from a nn pair to a $T = 0$ pn pair in $L_{NN} = 0$. For the excited states $^{10}\text{Be}(2_1^+1)$ and $^{10}\text{Be}(2_2^+1)$, two sets of $J^\pi = \{1^+, 2^+, 3^+\}$ for the spin-isospin partners appear in the $T = 0$ spectra, but the 2^+0 states are strongly mixed with each other because they almost degenerate energetically. $^{10}\text{Be}(2_1^+1)$ has a rotating core with $L_{\text{core}} = 2$ and it has strong transition strength to $^{10}\text{B}(1_2^+0, 2_{1,2}^+0, 3_2^+0)$, which almost degenerate because there is no spin-orbit splitting for the $T = 0$ pn pairs in $[L_{\text{core}} = 2, L_{NN} = 0, S_{NN} = 1]_{J=1,2,3}$. $^{10}\text{Be}(2_2^+1)$ with a rotating $S = 0$ nn pair in $L_{NN} = 2$ has dominant transition strength to $^{10}\text{B}(1_3^+0, 2_{1,2}^+0, 3_1^+0)$, which show large spin-orbit splitting of the $S = 1$ pn pairs in $L_{NN} = 2$. As a result of the spin-orbit splitting, the 3^+0 state partnered with $^{10}\text{Be}(2_2^+1)$ comes down to the ground state of ^{10}B . This assignment is consistent with the experimental data of the strong GT transition for $^{10}\text{B}(3_1^+0) \rightarrow ^{10}\text{Be}(2_2^+1)$ measured by charge exchange reactions [20]. Strictly speaking, it is in principle unable to definitely define L_{core} and L_{NN} for $N = Z = \text{odd}$

nuclei with deformed cores because core nucleons and valence nucleons are identical fermions and indistinguishable in fully microscopic wavefunctions of identical fermions. Nevertheless, the GT transitions from $N = Z + 2$ neighbors are observables and they enable us to classify the final states in $T = 0$ $N = Z = \text{odd}$ nuclei in terms of $T = 0$ pn pairs in connection with nn pairs in the initial states of $N = Z + 2$ nuclei.

In ^{14}N spectra, low-lying states are understood as spin-isospin partners of ^{14}C for NN hole pairs in the ^{16}O core. $^{14}\text{C}(0_1^+1)$ has the strong GT transition not to the lowest 1^+0 state but to the excited 1^+0 state $^{14}\text{N}(1_2^+0)$ because these states have NN hole pairs in the same $L_{NN} = 0$ orbit. Then, the GT transition occurs from the $S = 0$ nn hole pair to the $S = 1$ pn hole pair. The GT transitions from $^{14}\text{C}(2_1^+1)$ to $^{14}\text{N}(1_1^+0, 2_1^+0, 3_1^+0)$ show spin-isospin-flip features of NN hole pairs. Indeed, $^{14}\text{N}(1_1^+0, 2_1^+0, 3_1^+0)$ spectra show the spin-orbit splitting of the $S = 1$ pn hole pairs in $L_{NN} = 2$. Note that the ordering 1_1^+0 , 2_1^+0 , and 3_1^+0 is opposite to that of the particle-particle pair case because the spin-orbit mean potentials for hole states are repulsive.

Our assignments are consistent with the strong GT transition for $^{14}\text{C}(0_1^+1) \rightarrow ^{14}\text{N}(1_2^+0)$ experimentally measured by charge exchange reactions. Moreover, for the transitions from $^{14}\text{N}(1_1^+0)$, relatively strong GT transitions to $^{14}\text{C}(2_1^+1)$ and $^{14}\text{C}(2_2^+1)$ have been observed by charge exchange reactions [19]. They support significant $L_{NN} = 2$ component in $^{14}\text{N}(1_1^+0)$ consistently with the present assignment though quantitative reproduction of the $B(\text{GT})$ values is not satisfactory in the present calculation.

For the GT transition between the ground states of ^{14}C and ^{14}N , the experimental $B(\text{GT}; ^{14}\text{C}(0_1^+1) \rightarrow ^{14}\text{N}(1_1^+0))$ is anomalously small as known as a long life problem of ^{14}C . The suppression of the GT transition of $^{14}\text{C}(0_1^+1) \rightarrow ^{14}\text{N}(1_1^+0)$ is partially understood by the NN pair picture in LS -coupling scheme that $^{14}\text{N}(1_1^+0)$ is not the spin-isospin partner of the $^{14}\text{C}(0_1^+1)$ but that of $^{14}\text{C}(2_1^+1)$ because of the large spin-orbit splitting for the $S = 1$ pn hole pairs in $L_{NN} = 2$. It is different from the $A = 6$ and $A = 10$ systems, in which the lowest 1^+0 state is the spin-isospin partner of the ground state of the $N = Z + 2$ nucleus. The GT transition from the $[L_{NN} = 0, S_{NN} = 0]_{J=0}$ component in $^{14}\text{C}(0_1^+1)$ to the $[L_{NN} = 2, S_{NN} = 1]_{J=1}$ component in $^{14}\text{N}(1_1^+0)$ is forbidden because of the difference $\Delta L_{NN} = 2$ in spatial configurations. In other words, the GT transition is suppressed because of the LS -coupling pair correlation. Indeed, the calculated $B(\text{GT}) = 0.30$ is factor one smaller than the sum rule value and less than the half of the jj -coupling limit $B(\text{GT}) = 2/3$ for the pure $p_{1/2}^{-2}$ con-

figuration without the pair correlation. Our result for $B(\text{GT}; {}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0))$ is the same order as those of a NCSM calculation [23] and AMD+VAP calculation [24] but still largely overestimates the experimental data. In the present calculation, the NN pairs in ${}^{14}\text{C}(0_1^+1)$ and ${}^{14}\text{N}(1_1^+0)$ dominantly have $[L_{NN} = 0, S_{NN} = 0]_{J=0}$ and $[L_{NN} = 2, S_{NN} = 1]_{J=1}$ components, respectively, but they are not necessarily ideal LS -coupling pairs. Moreover, $[L_{NN} = 2, S_{NN} = 1]_{J=1}$ and $[L_{NN} = 0, S_{NN} = 1]_{J=1}$ are somewhat mixed with each other in the obtained ${}^{14}\text{N}(1_1^+0)$ and ${}^{14}\text{N}(1_2^+0)$. As a result of significant mixing of configurations, the calculated GT transition ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0)$ does not vanish. Additional scenarios are required to solve the long-life problem of ${}^{14}\text{C}(0_1^+1)$.

In the present analysis, we can understand low-energy spectra of $A = 6$, $A = 10$, and $A = 14$ nuclei from the LS -coupling NN pair picture and assign spin-isospin partners not only for the 0^+1 initial states but also the 2^+1 initial states as shown in Fig. 3. The spin-orbit splitting of the $J^\pi T = 1^+0, 2^+0, 3^+0$ states with $L_{NN} = 2$ coupled with the intrinsic spin $S = 1$ of the NN pair is essential in the spectra of $N = Z = \text{odd}$ nuclei. In the systematics of the spin-orbit splitting shown in Fig. 3, we can see that the splitting becomes large as A increases. It implies that the LS -coupling NN pairs feel the stronger spin-orbit mean potential in heavier systems.

V. SUMMARY AND OUTLOOK

We have studied the Gamow-Teller transitions from $N = Z + 2$ neighbors to $N = Z = \text{odd}$ nuclei in the p -shell region by using $T\beta\gamma$ -AMD+GCM. We have obtained that the strong GT transitions exhausting more than 50% of the sum rule for ${}^6\text{He}(0_1^+1) \rightarrow {}^6\text{Li}(1_1^+0)$, ${}^{10}\text{Be}(0_1^+1) \rightarrow {}^{10}\text{B}(1_1^+0)$, and ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_2^+0)$. We have also found the concentration of the GT strengths of the transitions from 2_1^+1 states, ${}^6\text{He}(2_1^+1) \rightarrow {}^6\text{Li}(1_2^+0, 2_1^+0, 3_1^+0)$, ${}^{10}\text{Be}(2_1^+1) \rightarrow {}^{10}\text{B}(1_2^+0, 2_1^+0, 2_2^+0, 3_2^+0)$, ${}^{10}\text{Be}(2_2^+1) \rightarrow {}^{10}\text{B}(1_3^+0, 2_1^+0, 2_2^+0, 3_1^+0)$, and ${}^{14}\text{C}(2_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0, 2_1^+0, 3_1^+0)$. These states connected with the strong GT transitions can be interpreted as “spin-isospin partner” states.

For further analysis, we have introduced two-nucleon-pair densities to visualize NN pair distributions, and found that $S = 0, T = 1$ nn pairs and $S = 1, T = 0$ pn pairs are dominantly formed in the $N = Z + 2$ and $N = Z = \text{odd}$ nuclei, respectively. We have studied the spin and orbital configurations of the NN pairs in LS -coupling scheme

TABLE III. Expectation values ($\langle S^2 \rangle$ and $\langle L^2 \rangle$) of the squared intrinsic spin and orbital angular momentum for ${}^6\text{Li}$, ${}^{10}\text{B}$, and ${}^{14}\text{N}$ obtained by $T\beta\gamma$ -AMD+GCM and ${}^6\text{He}$, ${}^{10}\text{Be}$, and ${}^{14}\text{C}$ obtained by $\beta\gamma$ -AMD+GCM.

$N = Z + 2$				$N = Z = \text{odd}$						
nuclide	$J_n^\pi T$	$\langle S^2 \rangle$	$\langle L^2 \rangle$	nuclide	$J_n^\pi T$	$\langle S^2 \rangle$	$\langle L^2 \rangle$	$J_n^\pi T$	$\langle S^2 \rangle$	$\langle L^2 \rangle$
${}^6\text{He}$	$0_1^+ 1$	0.12	0.12	${}^6\text{Li}$	$0_1^+ 1$	0.12	0.12	$1_1^+ 0$	1.97	0.06
	$2_1^+ 1$	0.19	5.65		$2_1^+ 1$	0.20	5.64	$1_2^+ 0$	1.90	5.75
								$2_1^+ 0$	2.00	5.99
								$3_1^+ 0$	2.01	6.01
${}^{10}\text{Be}$	$0_1^+ 1$	0.34	0.34	${}^{10}\text{B}$	$0_1^+ 1$	0.28	0.28	$1_1^+ 0$	1.94	0.35
	$2_1^+ 1$	0.30	6.00		$2_1^+ 1$	0.27	6.04	$1_2^+ 0$	1.92	5.43
								$2_1^+ 0$	2.02	6.49
								$3_2^+ 0$	1.97	7.53
	$2_2^+ 1$	0.12	6.11		$2_2^+ 1$	0.10	6.08	$1_3^+ 0$	1.99	5.94
								$2_2^+ 0$	2.02	6.61
								$3_1^+ 0$	2.05	7.15
${}^{14}\text{C}$	$0_1^+ 1$	0.55	0.55	${}^{14}\text{N}$	$0_1^+ 1$	0.61	0.61	$1_2^+ 0$	1.94	0.44
	$2_1^+ 1$	0.19	5.79		$2_1^+ 1$	0.21	5.83	$1_1^+ 0$	1.89	5.56
								$2_1^+ 0$	2.01	6.07
								$3_1^+ 0$	2.02	6.22

and discussed the behaviors of the LS -coupling NN pairs in relation to the GT transitions. The ground states of $N = Z + 2$ nuclei, ${}^6\text{He}(0_1^+ 1)$, ${}^{10}\text{Be}(0_1^+ 1)$, and ${}^{14}\text{C}(0_1^+ 1)$, and their partner states, ${}^6\text{Li}(1_1^+ 0)$, ${}^{10}\text{B}(1_1^+ 0)$, and ${}^{14}\text{N}(1_2^+ 0)$, have major $L = 0$ components, in which both the NN pairs and the core nuclei are in $L = 0$ states. The excited states, ${}^6\text{He}(2_1^+ 1)$, ${}^{10}\text{Be}(2_2^+ 1)$, and ${}^{14}\text{C}(2_1^+ 1)$, and their partner states have dominantly $L = 2$ components mainly contributed by the NN rotation around the core, whereas ${}^{10}\text{Be}(2_1^+ 1)$ and its spin-isospin partners have $L = 2$ components with the deformed 2α core rotating in $L = 2$. Based on the LS -coupling NN pairs, the strong GT transitions between spin-isospin partners can be understood as spin-isospin-flip phenomena from the $S = 0, T = 1$

nn pairs in the $N = Z + 2$ initial states to $S = 1, T = 0$ pn pairs in the $N = Z = \text{odd}$ final states. Namely, the transitions ${}^6\text{He}(0_1^+1) \rightarrow {}^6\text{Li}(1_1^+0)$, ${}^{10}\text{Be}(0_1^+1) \rightarrow {}^{10}\text{B}(1_1^+0)$ and ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_2^+0)$ are spin-flip phenomena of the NN pairs with $[L_{\text{core}} = 0, L_{NN} = 0]_{L=0}$, whereas ${}^6\text{He}(2_1^+1) \rightarrow {}^6\text{Li}(1_2^+0, 2_1^+0, 3_1^+0)$, ${}^{10}\text{Be}(2_2^+1) \rightarrow {}^{10}\text{B}(1_3^+0, 2_1^+0, 2_2^+0, 3_1^+0)$, and ${}^{14}\text{C}(2_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0, 2_1^+0, 3_1^+0)$ are those with $[L_{\text{core}} = 0, L_{NN} = 2]_{L=2}$. In the latter cases, the spectra of three final states with $J^\pi = 1^+, 2^+, 3^+$ are split because of the spin-orbit interaction for the $S = 1, T = 0$ pn pairs in $L_{NN} = 2$ wave. This spin-orbit splitting plays an important role in the low-energy spectra of the $N = Z$ odd nuclei. On the other hand, the spectra of ${}^{10}\text{B}(1_2^+0, 2_1^+0, 2_2^+0, 3_2^+0)$ partnered with ${}^{10}\text{Be}(2_1^+1)$ show small splitting because these states have dominant $[L_{\text{core}} = 2, L_{NN} = 0]_{L=2}$ component, in which the spin-orbit interaction does not affect the $S = 1, T = 0$ pn pairs in $L_{NN} = 0$ wave.

In comparison with experimental data, the magnetic moments μ and the magnetic dipole transition strengths $B(M1)$ are reasonably reproduced in the present calculation. Moreover, relatively enhanced $B(\text{GT})$ for ${}^6\text{He}(0_1^+1) \rightarrow {}^6\text{Li}(1_1^+0)$, ${}^{10}\text{Be}(0_1^+1) \rightarrow {}^{10}\text{B}(1_1^+0)$, and ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_2^+0)$ show consistent features with the present results. The present calculation also succeeds in describing the concentrations of the GT strengths from the $J = 2$ excited states: ${}^{10}\text{Be}(2_2^+1) \rightarrow {}^{10}\text{B}(3_1^+0)$ and ${}^{14}\text{C}(2_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0)$.

The present framework, $T\beta\gamma$ -AMD+GCM, is a useful tool to systematically study the pn pair correlations in $A = 6$, $A = 10$, and $A = 14$ nuclei. With this method, we can deal with nuclear deformations of the core nuclei and NN pair formation in the same footing. This is one of the great advantages superior to three-body models with a spherical inert core.

As mentioned above, the strong GT transitions can be understood in terms of the transitions $nn(S = 0, T = 1) \rightarrow pn(S = 1, T = 0)$ in LS -coupling scheme. It means that the GT transitions is a good probe to clarify the dynamics of the pn pairs in $N = Z = \text{odd}$ nuclei through the connection with the nn pairs in the neighboring nuclei. We should comment that anomalous suppression of the GT transition ${}^{14}\text{C}(0_1^+1) \rightarrow {}^{14}\text{N}(1_1^+0)$ is not reproduced in the present calculation and it is still a remaining problem.

In light mass nuclei, the LS -coupling pn pairs are formed. However, for heavier nuclei, the description of pn pair correlation in LS -coupling scheme is no longer valid because jj -coupling pn pairs and also the pn pair condensation are expected because of the spin-orbit interactions. Further investigations of $N = Z = \text{odd}$ nuclei in a wide mass number region from light to heavy mass nuclei are required for deeper understanding of pn pair correlations.

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